Spin-Dependent Transport of Electrons in a Shuttle Structure

L. Y. Gorelik,* R. I. Shekhter, and M. Jonson

Department of Applied Physics, Chalmers University of Technology and Göteborg University, SE-412 96 Göteborg, Sweden

S. I. Kulinich

Department of Applied Physics, Chalmers University of Technology and Göteborg University, SE-412 96 Göteborg, Sweden and B. I. Verkin Institute for Low Temperature Physics and Engineering, 47 Lenin Avenue, 61103 Kharkov, Ukraine

V. M. Vinokur

Materials Science Division, Argonne National Laboratory, 9700 South Cass, Argonne, Illinois 6043

(Dated: February 2, 2008)

We consider "shuttling" of spin-polarized electrons between two magnetic electrodes (half-metals) by a movable dot with a single electronic level. If the magnetization of the electrodes is antiparallel we show that the transmittance of the system can be changed by orders of magnitude if an external magnetic field, perpendicular to the polarization of the electronic spins, is applied. A giant magnetotransmittance effect can be achieved for weak external fields of order $1 \div 10$ Oe.

PACS numbers: 72.25.Hg, 73.43.Jn., 73.61.Ey, 72.50.Bb

INTRODUCTION

Metal-organic nanocomposite materials are interesting from the point of view of the "bottom-up" approach to building future electronic devices. The ability of the organic parts of the composite materials to identify and latch on to other organic molecules is the basis for the possible self assembly of nanoscale devices, while the metallic components provide mechanical robustness and improve the electrical conductance.

Such composite materials are heteroelastic in the sense that the mechanical rigidity of

the organic and metallic components are very different. This allows for a special type of deformation, where hard metallic components embedded in a soft organic matrix can be rearranged in space at a low deformation energy cost associated with stretching and compressing the soft matrix. Strong Coulomb forces, due to accumulation of electronic charge in embedded nanoscale metallic particles, can be a source of such mechanical deformations. This leads to a scenario where the transport of electric charge, possibly due to tunnelling of electrons between metal particles, becomes a complex nano-electromechanical phenomenon, involving an interplay of electronic and mechanical degrees of freedom [1]. Such an interplay can lead to new physics, as was recently demonstrated theoretically for the simplest possible structure — a Nanoelectromechanical single-electron transistor. The electromechanical instability predicted to occur in this device at large enough bias voltage was shown to provide a new mechanism of charge transport [2]. This mechanism can be viewed as a "shuttling" of single electrons by a metallic island — a Coulomb dot — suspended between two metal electrodes. The predicted instabilily leads to a periodic motion of the island between the electrodes shuttling charge from one to the other.

The shuttle instability appears to be a rather general phenomenon. It has, *e.g.*, been shown to occur even for extremely small suspended metallic particles (or molecules) for which the coherent quantum dynamics of the tunnelling electrons [3] or even the quantum dynamics of the mechanical vibration [4, 5, 6, 7] become essential. Nanomechanical transport of electronic charge can, however, occur without any such instability, *e.g.*, in an externally driven device containing a cantilever vibrating at frequencies of order 100 MHz. A small metallic island attached to the tip of the vibrating cantilever may shuttle electrons between metallic leads as has recently been demonstrated [8]. Further experiments with magnetic and superconducting externally driven shuttles as suggested in [9], seem to be a natural extension of this work. Fullerene-based nanomechanical structures [10] are also of considerable interest.

The possibility to place transition-metal atoms or ions inside organic molecules introduces a new "magnetic" degree of freedom that allows the electronic spins to be coupled to mechanical and charge degrees of freedom [12]. By manipulating the interaction between the spin and external magnetic fields and/or the internal interaction in magnetic materials, spin-controlled nanoelectromechanics may be achieved. An inverse phenomenon — nanomechanical manipulation of nanomagnets — was suggested earlier in [11]. A magnetic field, by inducing the spin of electrons to rotate (precess) at a certain frequency, provides a clock

for studying the shuttle dynamics and a basis for a dc spectroscopy of the corresponding nanomechanical vibrations.

A particularly interesting situation arises when electrons are shuttled between electrodes that are half-metals. A half-metal is a material that not only has a net magnetization as do ferromagnets, but all the electrons are in the same spin state — the material is fully spin-polarized. Examples of such materials can be found among the perovskite maganese oxides, a class of materials that show an intrinsic, so called "colossal magnetoresistance" effect at high magnetic fields (of order 10-100 kOe) [13].

A large magnetoresistance effect at lower magnetic fields has been observed in layered tunnel structures where two thin perovskite manganese oxide films are separated by a tunnel barrier [13, 14, 15, 16, 17]. Here the spin polarization of electronic states crucially affects the tunnelling between the magnetic electrodes. This is because electrons that can be extracted from the source electrode have there spins aligned in a definite direction, while electrons that can be injected into the drain electrode must also have there spins aligned — possibly in a different direction. Clearly the tunnelling probability and hence the resistance must be strongly dependent on the relative orientation of the magnetization of the two electrodes. An external magnetic field aligns the magnetization direction of the two films at different field strengths, so that the relative magnetization can be changed between high- and low resistance configurations. A change in the resistance of trilayer devices by factors of order 2-5 have in this way been induced by magnetic fields of order 200 Oe [14, 15, 16]. The required field strength is determined by the coercivities of the magnetic layers. This makes it difficult to use a tunneling device of the described type for sensing very low magnetic fields. In this paper we propose a new functional principle — spin-dependent shuttling of electrons — for low-magnetic field sensing purposes. We will show that this principle can lead to a giant magnetoresistance effect in external fields as low as 1-10 Oe.

The new idea which we propose to pursue is to use the external magnetic field to manipulate the *spin of shuttled electrons* rather than the magnetization of the leads. The possibility to "trap" electrons on a nanomechanical shuttle (decoupled from the magnetic leads) during quite a long time on the scale of the time it takes an electron to tunnel on/off the shuttle makes it possible for even a weak external field to rotate the electron's spin to a significant degree. Such a rotation allows the spin of an electron, loaded onto the shuttle from the spin-polarized source electrode, to be reoriented in order to allow the electron finally to tunnel

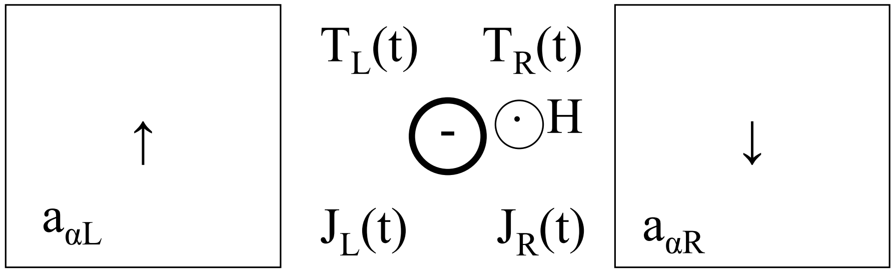

FIG. 1: Schematic view of the nanomechanical GMR device: a movable dot with a single electron level couples to the leads due to tunnelling of electrons, described by the tunnelling probability amplitudes $T_{L,R}(t)$), and due to the exchange interaction whose strength is denoted by $J_{L,R}(t)$. An external magnetic field H is oriented perpendicular to the direction of the magnetization in the leads (arrows).

from the shuttle to the spin-polarized drain lead. As we will show below, the magnetic field induced spin-rotation of shuttled electrons is a very sensitive nanomechanical mechanism for a giant magnetoresistance (GMR) effect.

FORMULATION OF THE PROBLEM. GENERAL EXPRESSION FOR THE CURRENT

A schematic view of the nanomechanical GMR device to be considered is presented in Fig. 1. Two fully spin-polarized magnets with fully spin-polarized electrons serve as source and drain electrodes in a tunnelling device. In this paper we will consider the situation when the electrodes have exactly opposite polarization. A mechanically movable quantum dot (described by a time-dependent displacement $x(t)$), where a single energy level is available for electrons, performs forced harmonic oscillations with period $T = 2\pi/\omega$ between the leads. The external magnetic field is perpendicular to the orientation of the magnetization in both leads.

The Hamiltonian that governs the dynamical evolution of the system is

$$\begin{aligned}\hat{\mathcal{H}}(t) &= \varepsilon_0(a^\dagger_\uparrow a_\uparrow + a^\dagger_\downarrow a_\downarrow) + \sum_\alpha(\varepsilon_\alpha a^\dagger_{\alpha,L}a_{\alpha,L} + \varepsilon_\alpha a^\dagger_{\alpha,R}a_{\alpha,R}) \\ &- J_L(t)(a^\dagger_\uparrow a_\uparrow - a^\dagger_\downarrow a_\downarrow) - J_R(t)(a^\dagger_\downarrow a_\downarrow - a^\dagger_\uparrow a_\uparrow) - (g\mu H/2)(a^\dagger_\uparrow a_\downarrow + a^\dagger_\downarrow a_\uparrow) \\ &+ T_L(t)\sum_\alpha(a^\dagger_{\alpha,L}a_\uparrow + a^\dagger_\uparrow a_{\alpha,L}) + T_R(t)\sum_\alpha(a^\dagger_{\alpha,R}a_\downarrow + a^\dagger_\downarrow a_{\alpha,R})\,,\end{aligned} \tag{1}$$

where $a^\dagger_{\alpha,L(R)}$, $(a_{\alpha,L(R)})$ are the creation (annihilation) operators of electrons with the energy ε_α on the left (right) lead (we have suppressed the spin indices for the electronic states in the leads due to the assumption of full spin polarization), $a^\dagger_{\uparrow(\downarrow)}(a_{\uparrow(\downarrow)})$ are the creation (annihilation) operators on the dot, ε_0 is the energy of the on-dot level, $J_{L(R)}(t) \equiv J_{L(R)}(x(t))$ are the exchange interactions between the on-grain electron and the left (right) lead, $\Lambda_{L(R)}(t) \equiv \Lambda_{L(R)}(x(t))$ are the tunnel coupling amplitudes, g is the gyromagnetic ratio and μ is the Bohr magneton.

The single-electron density matrix describing electronic transport between the leads may be presented in the form:

$$\hat{\rho} = \sum_\alpha w_{\alpha,L}|\Psi^{\alpha,L}\rangle\langle\Psi^{\alpha,L}| + \sum_\alpha w_{\alpha,R}|\Psi^{\alpha,R}\rangle\langle\Psi^{\alpha,R}|\,. \tag{2}$$

Here $|\Psi^{\alpha,L}\rangle$ are single-electron states that obey the time-dependent Shrödinger ($\hbar = 1$) equation with a Hamiltonian given by Eq. (1). The initial condition has the form

$$|\Psi^{\alpha,L(R)}(t \to -\infty)\rangle = |\alpha, L(R)\rangle \exp(-i\varepsilon_\alpha t)\,,$$

where $|\alpha, L(R)\rangle$ is a single-electron state on the left (right) lead with energy ε_α.

We will suppose that the internal relaxation in the leads is fast enough to lead to equilibrium distributions of the electrons. This means that $w_{\alpha,L(R)} = f(\varepsilon_\alpha \mp V/2)$ (where $f(\varepsilon)$ is the Fermi distribution function) and V is applied voltage.

The problem at hand is greatly simplified if one considers the large bias-voltage limit

$$\mid V - \varepsilon_0 \mid \gg \nu\Lambda^2_{\max}\,, \tag{3}$$

where ν is the density of states on the leads. The restriction(3) does not allow us to consider a narrow transition region of voltages from the zero-current regime at $V < \varepsilon_0 - \nu\Lambda^2_{\max}$ to the fully transmissive (in the absence of spin polarization effects) regime at $V > \varepsilon_0 + \nu\Lambda^2_{\max}$. However, it covers the in practise most important case when the fully transmissive junction

is strongly affected by electronic spin-polarization. Therefore, in our further considerations we will take $w_{\alpha,L} = 1$, $w_{\alpha,R} = 0$ and $\varepsilon_0 = 0$.

We will calculate the average current, I, through the system from the relation

$$I = \frac{1}{T}\int_0^T dt\, \mathrm{Tr}\left\{\hat{\rho}\hat{\jmath}\right\}, \tag{4}$$

$$\hat{\jmath} = e\frac{\partial \hat{N}_R}{\partial t} = ie[\hat{\mathcal{H}}, \hat{N}_R] = ieT_R(t)\sum_\alpha (a_\downarrow^\dagger a_{\alpha,R} - a_{\alpha,R}^\dagger a_\downarrow)\,,$$

where $\hat{N}$ is the electron number operator for the right lead, $\hat{N} = \sum_\alpha a_{\alpha,R}^\dagger a_{\alpha,R}$.

In general, the state $|\Psi^{\alpha,L}\rangle$ can be expressed as

$$|\Psi^{\alpha,L}(t)\rangle = c_\uparrow^\alpha(t)|\uparrow\rangle + c_\downarrow^\alpha(t)|\downarrow\rangle + \sum_\beta (c_L^{\alpha,\beta}(t)|\beta,L\rangle + c_R^{\alpha,\beta}(t)|\beta,R\rangle)\,, \tag{5}$$

Thus the problem is reduced to determining the coefficients $c_{R(L)}^{\alpha,\beta}$ and $c_{\downarrow(\uparrow)}^\alpha$.

At this point it is convenient to introduce the bi-vectors

$$\mathbf{c}^\alpha = \begin{pmatrix} c_\uparrow^\alpha \\ c_\downarrow^\alpha \end{pmatrix}, \quad \mathbf{e}_1 = \begin{pmatrix} 1 \\ 0 \end{pmatrix}, \quad \text{and} \quad \mathbf{e}_2 = \begin{pmatrix} 0 \\ 1 \end{pmatrix},$$

so that the coefficients $c_{R(L)}^{\alpha,\beta}$ can be expressed as (see Appendix 1)

$$c_L^{\alpha,\beta} = e^{-i\varepsilon_\beta t}\delta_{\alpha\beta} - i\int_{-\infty}^t dt' e^{i\varepsilon_\beta(t-t')}T_L(t')(\mathbf{e}_1, \mathbf{c}^\alpha(t'))\,,$$
$$c_R^{\alpha,\beta} = -i\int_{-\infty}^t dt' e^{i\varepsilon_\beta(t-t')}T_R(t')(\mathbf{e}_2, \mathbf{c}^\alpha(t'))\,.$$

Here $(\mathbf{a}, \mathbf{b})$ is the inner product of two bi-vectors. As shown in Appendix 1, by using the wide band approximation (*i.e.* by taking the electron density of states in the leads ν to be constant) the equation for the bi-vectors $\mathbf{c}^\alpha$ takes the form

$$i\frac{\partial \mathbf{c}^\alpha}{\partial t} = \hat{R}(t)\mathbf{c}^\alpha + \mathbf{f}^\alpha(t)\,. \tag{6}$$

Here $\mathbf{f}^\alpha(t) = T_L(t)e^{-i\varepsilon_\alpha t}\mathbf{e}_1$ and the matrix $\hat{R}(t)$ is

$$\hat{R}(t) = \begin{pmatrix} -J(t) - i\Gamma_L(t)/2 & -g\mu H/2 \\ -g\mu H/2 & J(t) - i\Gamma_R(t)/2 \end{pmatrix}, \tag{7}$$

where $J(t) = J_L(t) - J_R(t)$ and $\Gamma_{L(R)}(t) = 2\pi\nu\Lambda_{L(R)}^2(t)$ is the level width.

The formal solution of Eq. (6) can be written in the form

$$\mathbf{c}^{\alpha}(t) = -i\int_{-\infty}^{t} dt'\hat{L}(t,t')\,\mathbf{f}^{\alpha}(t')\,, \tag{8}$$

where the "evolution" operator $\hat{L}(t,t')$, ($\hat{L}(t,t)=\hat{I}$), is defined as the solution of the equation

$$i\frac{\partial \hat{L}(t,t')}{\partial t} = \hat{R}(t)\hat{L}(t,t')\,, \tag{9}$$

and obeys the multiplicative and periodicity properties,

$$\hat{L}(t,t') = \hat{L}(t,t'')\hat{L}(t'',t'), \quad \hat{L}(t+T,t'+T) = \hat{L}(t,t'). \tag{10}$$

Using Eq. (8) together with Eq. (4), one can write the average current on the form

$$I = \frac{e}{T}\int_{0}^{T} dt\Gamma_R(t)\int_{-\infty}^{t} dt'\Gamma_L(t')|\hat{L}_{21}(t,t')|^2\,, \tag{11}$$

where $\hat{L}_{21}(t,t')$ is a matrix element of the operator $\hat{L}(t,t')$; $\hat{L}_{21}(t,t') = (\mathbf{e}_2, \hat{L}(t,t')\mathbf{e}_1)$.

Since the probability amplitude for tunnelling is exponentially sensitive to the position of the dot, the maximum of the tunnel exchange interaction between an electron on the dot and an electron in one lead occurs when the tunnelling coupling to the other lead is negligible. This is why we will assume the following property of tunnelling amplitude $\Lambda_{L,R}(t)$ to be fulfilled:

$$T_L(t)T_R(t) = 0,\ T_L(t), T_R(t) \neq 0 \tag{12}$$

This assumption allows us to divide the time interval $(0,T)$ into the intervals $(0,\tau) + (\tau, T/2) + (T/2, T/2+\tau) + (T/2+\tau, T)$. We suppose that $T_L(t)\neq 0$ (but $H=0$) only in the time interval $(0,\tau)$ (and, analogously, $T_R(t)\equiv T_L(t+T/2)\neq 0$ in the time interval $(T/2, T/2+\tau)$). Using this approximation together with the properties (10) of the operator $\hat{L}(t,t')$, we arrive at the following expression for the average current (Appendix 2):

$$I = \frac{e}{T}(1-e^{-\Gamma})^2\sum_{n=0}^{\infty}|(\mathbf{e}_2, \hat{L}(T/2,\tau)\hat{L}^n\mathbf{e}_1)|^2. \tag{13}$$

Here $\hat{L}\equiv\hat{L}(T+\tau,\tau)$ and

$$\Gamma = 2\pi\nu\int_0^{\tau} dt T_L^2(t) \tag{14}$$

is the tunnelling rate. Consequently, in order the calculate the average current it is necessary to investigate the properties of the evolution operator $\hat{L}$. It follows from its definition that

$$\begin{aligned} \hat{L} &= \hat{L}(T+\tau,\tau) = \hat{L}(T+\tau,T)\hat{L}(T,T/2+\tau)\hat{L}(T/2+\tau,T/2)\hat{L}(T/2,\tau) \\ &= e^{-(1+\sigma_3)\Gamma/4+i\sigma_3\Phi_0}\hat{L}(T,T/2+\tau)e^{-(1-\sigma_3)\Gamma/4-i\sigma_3\Phi_0}\hat{L}(T/2,\tau)\,, \end{aligned} \tag{15}$$

where $\Phi_0 = \int_0^\tau dt J(t)$. From the symmetry properties of the operator $\hat{R}(T/2+\tau<t<T)$,

$$\hat{R}^\dagger = \hat{R}, \sigma_2 \hat{R}^* = -\hat{R}, \sigma_3 \hat{R}(-t) = -\hat{R}(t)\sigma_3$$

it follows that the operator $\hat{U} \equiv \hat{L}(T, T/2+\tau)$ has the form

$$\hat{U} = \begin{pmatrix} \sqrt{1-\gamma^2} & i\gamma e^{i\varphi} \\ i\gamma e^{-i\varphi} & \sqrt{1-\gamma^2} \end{pmatrix} \tag{16}$$

In addition to this, $\hat{L}(T/2,\tau) = \sigma_1 \hat{U} \sigma_1$. As a result, the operator $\hat{L}$ can be expressed as

$$\hat{L} = e^{-\Gamma/2} \left(e^{-\sigma_3 \Gamma/4 + i\Phi_0\sigma_3} \hat{U} \sigma_1 \right)^2 . \tag{17}$$

Proceeding with the analysis we (i) calculate the eigenvalues λ_i and eigenvectors $\mathbf{b}_i$ of the operator $\hat{L}$ of Eq. (17); $\hat{L}\mathbf{b}_i = \lambda_i \mathbf{b}_i$, (ii) substitute the expansion $\mathbf{e}_i = a_{ji}\mathbf{b}_j$ (where $(a)^{-1} = (\mathbf{e}_i, \mathbf{b}_j)$) into Eq. (13) and calculate the average current. The result is

$$I = \frac{e\kappa}{T} \sinh\Gamma/2 \frac{\cosh\Gamma/2 + \cos 2\vartheta}{\sinh^2\Gamma/2 + \kappa(1+\cos 2\vartheta \cosh\Gamma/2)} , \tag{18}$$

where $\vartheta = \varphi + \Phi_0, \kappa = 2\gamma^2/(1+\gamma^2)$. Equation (18) for the average current is our main result.

CALCULATION OF THE CURRENT IN LIMIT OF STRONG AND WEAK EXCHANGE COUPLING BETWEEN THE DOT AND THE LEADS

Although the result (18) for the tunnel current is both transparent and compact, it is in general a rather complicated problem to find the magnetic field dependence of the coefficient κ, which depends on the probability amplitude γ for flipping the spin of shuttled electrons. Three different time scales are involved in the spin dynamics of a shuttled electron. They correspond to three characteristic frequencies: (i) the frequency of spin rotation, determined by the tunnel exchange interaction with the magnetic leads; (ii) the frequency of spin rotation in the external magnetic field, and (iii) the frequency of shuttle vibrations. Different regimes occur depending on the relation between these time scales. Here we will consider two limiting cases, where a simple solution of the problem can be found. Those are the limits of weak $J_{L(R)} \ll \mu H$ and strong $J_{L(R)} \gg \mu H$ exchange interactions with the leads.

Weak exchange interaction

In the limit $J_{L(R)} \ll \mu H$ one may neglect the influence of the magnetic leads on the on-dot electron spin dynamics. In this case the matrix $\hat{U}$ given by Eq. (16) can easily be calculated and Eq. (18) reduces to

$$I = \frac{2e}{T} \frac{\sin^2 \vartheta/2 \, \tanh \Gamma/4}{\sin^2 \vartheta/2 + \tanh^2 \Gamma/4}\,, \tag{19}$$

where $\vartheta = g\mu \int_\tau^{T/2} dt H$ is the rotation angle of the spin in the external field.

Two different scales for the external magnetic field determine the magneto-transmittance in this limit. One scale is associated with the width of the resonant magnetic field dependence (see the denominator in Eq. (19)). This scale is (restoring dimension)

$$\delta H = \Gamma \frac{\hbar\omega}{g\mu}\,, \tag{20}$$

where ω is the shuttle vibration frequency. The second scale,

$$\Delta H = \frac{\hbar\omega}{g\mu}\,, \tag{21}$$

comes from the periodic function $\sin^2 \vartheta/2$ that enters Eq. (19). The magnetic-field dependence of the current is presented in Fig. 2a. Dips in the transmittance of width δH appear periodically as the magnetic field is varied, the period being Δ. This amount to a giant magneto-transmittance effect. It is interesting to notice that by measuring the period of the variations in $I(H)$ one can in principle determine the shuttle vibration frequency. This amounts to a dc method for spectroscopy of the nanomechanical vibrations. Equation (21) gives a simple relation between the vibration frequency and the period of the current variations. The physical meaning of this relation is very simple: every time when $\omega/\Omega = n + 1/2$ (Ω is the spin precession frequency in a magnetic field) the shuttled electron is able to fully flip its spin to remove the "spin-blockade" of tunnelling between spin polarized leads having their magnetization in opposite directions.

Strong Dot-Leads Exchange Interaction

A strong magnetic coupling to the leads, $J_{\max} \gg \mu H$, preserves the electron spin polarization, preventing spin-flips of shuttled electrons due to an external magnetic field. However,

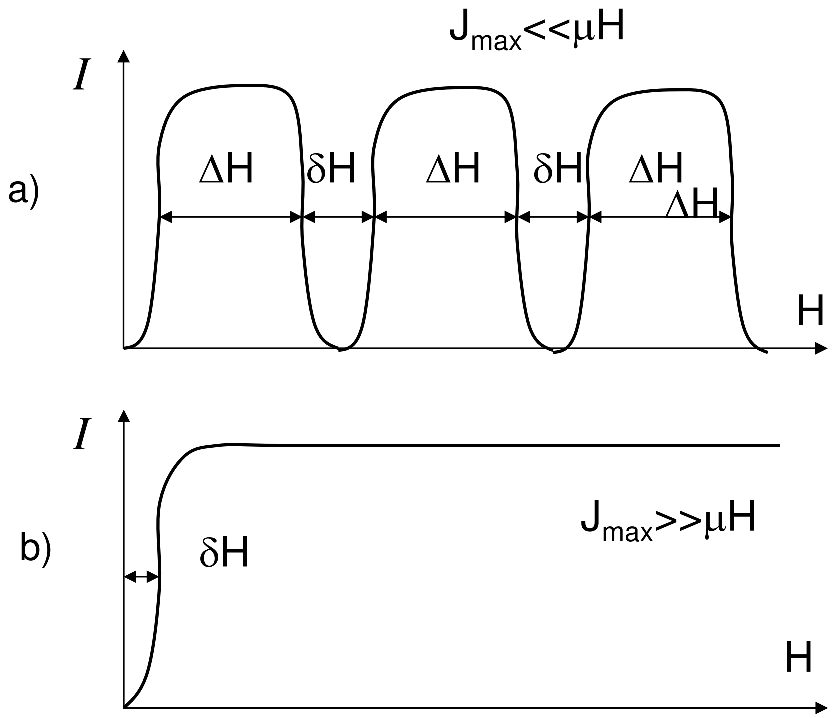

FIG. 2: Magnetic-field dependence of the transmittance of the device shown in Fig. 1 for the limiting cases of a) weak and b) strong exchange coupling between dot and leads. The period ΔH and the width δH of the "dips" are given by Eqs. (21) and (20) for case a) and δH is given by Eq. (24) for the case b).

if the magnetization of the two leads are in opposite directions, the exchange coupling to the leads have different sign. Therefore, the exchange couplings to the two leads tend to cancel out when the dot is in the middle of the junction. Hence the strong exchange interaction affecting a dot electron depends on time and periodically changes sign, being arbitrary small close to the time of sign reversal. In Fig. 3 the on-dot electronic energy levels for spins parallel and antiparallel to the lead magnetization are presented as a function of time. The effect of an external magnetic field is in the limit $J_{L(R)} \gg \mu H$ negligible almost everywhere, except in the vicinity of the level crossing. At this "time point" ,which we denote t_{LZ}, the external magnetic field removes the degeneracy and a gap is formed in the spectrum (dashed curve). The probability of electronic spin-flip in this case is determined by the probabil-

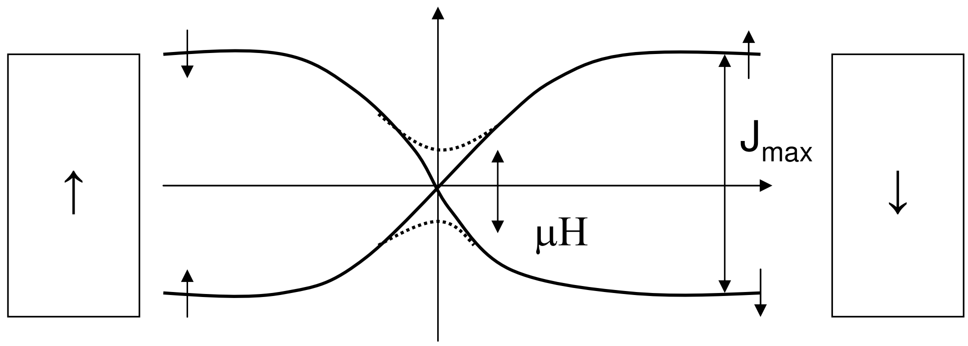

FIG. 3: On-dot energy levels for spin-up and spin-down electron states as a function of the position of the dot. Level crossing in the middle of the device is removed by an external magnetic field.

ity of a Landau-Zener reflection from the gap formed by the magnetic field (in this case a Landau-Zener transition across the gap is a mechanism for backscattering of the electron, since this is the channel where the electronic spin is preserved). The matrix $\hat{U}$ can readily be expressed in terms of Landau-Zener scattering amplitudes. The amplitude and phase of electronic spin-flip is given by $\varphi = \varphi_0 + \Phi_1, \varphi_0$ is the Landau-Zener phase shift,

$$\Phi_1 = \int_{\tau}^{T/2-\tau} dt J(t) \tag{22}$$

and γ^2 is the probability of the Landau-Zener "backward" scattering,

$$\gamma^2 = 1 - \exp\left[-\frac{\pi(\mu H)^2}{J'(t_{LZ})}\right] . \tag{23}$$

Schematical view of $I(H)$ dependence is presented on a Fig.2b. The width δH of the minimum in $I(H)$ dependence can be found directly from Eqs.(18), (23)

$$\delta H = \frac{\pi g \mu}{\sqrt{J_0 \hbar\omega}} , \tag{24}$$

where $J_0 = \min(J_{L(R)}(t))$.

CONCLUSION

The analysis presented above demonstrates the possibility of a giant magneto-transmittance effect caused by shuttling of spin-polarized electrons between magnetic source-

and drain electrodes. The sensitivity of the shuttle current to an external magnetic field is determined, according to Eq. (20), by the transparency of the tunnel barriers. By diminishing the tunnelling transmittance one can increase the sensitivity of the device to an external magnetic field. The necessity to have a measurable current determines the limit of this sensitivity. In the low transparency limit, $\Gamma \ll 1$, the current through the device can be estimated as $I \simeq e\Gamma\omega$. If one denotes the critical field that determines the sensitivity of the device by H_{cr}, one finds from Eq. (20) that $H_{cr} \simeq \delta H$. The critical field can now be expressed in terms of the current transmitted through the device as

$$H_{cr}(\mathrm{Oe}) \simeq \frac{\hbar I}{e\mu g} \simeq \frac{g_0}{g}(3\times 10^2) I(\mathrm{n\AA})\,, \tag{25}$$

where $g_0\,(=2)$ is the gyromagnetic ratio for the free electrons. For $I \simeq 10^{-1} \div 10^{-2}$ nA and $g_0/g \simeq 1/3$ this gives a range $H_{cr} \simeq 1 \div 10\,\mathrm{Oe}$. A further increase in sensitivity would follow if one could use a shuttle with several (N) electronic levels involved in the tunnelling process. The critical magnetic field would then be inversely proportional to the number of levels, $H_{cr}(N) = H_{cr}(N=1)/N$.

Acknowledgment

Financial support from the Royal Swedish Academy of sciences (SIK), the Swedish Science Research Council (LIG, RIS) and the Swedish Foundation for Strategic Research (LIG, RIS, MJ) is gratefully acknowledged. Support from the U.S.Department of energy (Basic Energy Science, Contract No W-31-109-ENG-38) is gratefully acknowledged.

Appendix 1

The Shrödinger equation results in equations for the coefficients $c^{\alpha\beta}_{R(L)}, c^{\alpha}_{\uparrow(\downarrow)}$:

$$\begin{aligned}
i\frac{\partial c^\alpha_\uparrow}{\partial t} &= -J(t)c^\alpha_\uparrow - (g\mu H/2)c^\alpha_\downarrow + T_L(t)\sum_\beta c^{\alpha\beta}_L(t)\,, \\
i\frac{\partial c^\alpha_\downarrow}{\partial t} &= J(t)c^\alpha_\uparrow - (g\mu H/2)c^\alpha_\downarrow + T_R(t)\sum_\beta c^{\alpha\beta}_R(t)\,, \\
i\frac{\partial c^{\alpha\beta}_L}{\partial t} &= \varepsilon_\beta c^{\alpha\beta}_L + T_L(t)c^\alpha_\uparrow(t)\,, \\
i\frac{\partial c^{\alpha\beta}_R}{\partial t} &= \varepsilon_\beta c^{\alpha\beta}_R + T_R(t)c^\alpha_\downarrow(t)\,.
\end{aligned} \tag{26}$$

As it follows from the last two equations (together with the initial conditions)

$$c_L^{\alpha\beta}(t) = e^{-i\varepsilon_\beta t}\delta_{\alpha\beta} - i\int_{-\infty}^{t} dt' e^{i\varepsilon_\beta (t'-t)} T_L(t') c_\uparrow^\alpha(t') \,, \tag{27}$$

$$c_R^{\alpha\beta}(t) = -i\int_{-\infty}^{t} dt' e^{i\varepsilon_\beta (t'-t)} T_R(t') c_\downarrow^\alpha(t') \,.$$

Therefore, for the $\sum_\beta c_R^{\alpha\beta}(t)$ one gets

$$\sum_\beta c_R^{\alpha\beta}(t) = -i\int_{-\infty}^{t} dt' T_R(t') c_\downarrow^\alpha(t') \sum_\beta e^{i\varepsilon_\beta (t'-t)} \,.$$

In wide-band approximation we suppose $\nu(\varepsilon) = \text{const}$, therefore $\sum_\beta e^{i\varepsilon_\beta (t'-t)} = 2\pi\nu\delta(t'-t)$ and

$$\sum_\beta c_R^{\alpha\beta}(t) = -i\pi\nu T_R(t) c_\downarrow^\alpha \,. \tag{28}$$

Analogously,

$$\sum_\beta c_L^{\alpha\beta}(t) = e^{-i\varepsilon_\alpha t} - i\pi\nu T_L(t) c_\uparrow^\alpha \,. \tag{29}$$

Substitute the expressions, Eqs.(28), (29), to the first two equations (26), one get the equation Eq. (6) for the bi-vector $\mathbf{c}^\alpha$.

Appendix 2

Under our approximation we can change the integration limits in Eq. (11):

$$I = (2\pi\nu)^2 \frac{e}{T}\int_0^T dt T_R^2(t) \int_{-\infty}^{t} dt' T_L^2(t') |\hat{L}_{21}(t,t')|^2$$

$$= (2\pi\nu)^2 \frac{e}{T}\int_{T/2}^{T/2+\tau} dt T_R^2(t) \int_{-\infty}^{\tau} dt' T_L^2(t') |\hat{L}_{21}(t,t')|^2 \,. \tag{30}$$

Beside this, in the time moments $T/2 < t < T/2+\tau$ $\hat{L}(t,T/2)$ is a diagonal matrix. Therefore $\hat{L}_{21}(t,t') = \hat{L}_{22}(t,T/2)\hat{L}_{21}(T/2,t')$. As a consequence, the integral in the expression for the average current, Eq. (30), is factorized:

$$I = (2\pi\nu)^2 \frac{e}{T}\int_{T/2}^{T/2+\tau} dt T_R^2(t) |\hat{L}_{22}(t,T/2)|^2 \int_{-\infty}^{\tau} dt' T_L^2(t') |\hat{L}_{21}(T/2,t')|^2 \,. \tag{31}$$

The first integral in Eq. (31) is easy to calculate. Having in mind that $(T/2 < t < T/2+\tau)$

$$|\hat{L}_{22}(t,T/2)|^2 = \exp\left[-2\pi\nu \int_{T/2}^{t} dt T_R^2(t)\right] \,,$$

one gets

$$2\pi\nu\int_{T/2}^{T/2+\tau} dt T_R^2(t)|\hat{L}_{22}(t,T/2)|^2 = 1 - e^{-\Gamma}\,, \tag{32}$$

where quantity Γ is defined in Eq. (14).

The calculation of the second integral in Eq. (31) can be done in the same manner. One has the set of equalities,

$$\int_{-\infty}^{\tau} T_L^2(t)|\hat{L}_{21}(T/2,t)|^2 = \sum_{n=0}^{\infty}\int_{-nT}^{-nT+\tau} dt T_L^2(t)|\hat{L}_{21}(T/2,t)|^2$$
$$= \sum_{n=0}^{\infty}\int_0^{\tau} dt T_L^2(t)|(\mathbf{e}_2, \hat{L}(T/2,\tau)\hat{L}(\tau,t-nT)\mathbf{e}_1)|^2\,. \tag{33}$$

For the quantity $\hat{L}(\tau,t-nT) = \hat{L}(\tau+nT,t)$ one has

$$\hat{L}(\tau{+}nT,t) = \hat{L}(\tau{+}nT,\tau{+}(n{-}1)T)\hat{L}(\tau{+}(n{-}1)T,\tau{+}(n{-}2)T)...\hat{L}(\tau,t) = \hat{L}^n(\tau{+}T,\tau)\hat{L}(\tau,t)\,. \tag{34}$$

Therefore,

$$\int_{-\infty}^{\tau} dt T_L^2(t)|\hat{L}_{21}(T/2,t)|^2 = \sum_{n=0}^{\infty}\int_0^{\tau} dt T_L^2(t)|(\mathbf{e}_2, \hat{L}(T/2,\tau)\hat{L}^n\hat{L}(\tau,t)\mathbf{e}_1)|^2\,. \tag{35}$$

Calculating the integral in the same manner, as in Eq. (32), one gets the Eq. (13) for the average current.

* Electronic address: gorelik@fy.chalmers.se

[1] R. I. Shekhter, Yu. Galperin, L. Y. Gorelik et al., J.Phys.:Condens. Matter **15** R441(2003).

[2] L. Y. Gorelik, A. Isacsson, M. V. Voinova et al., Phys. Rev. Lett. **80**, 4526 (1998).

[3] D. Fedorets, L. Y. Gorelik, R. I. Shekhter, and M. Jonson, Europhys. Lett. **58**, 99 (2002).

[4] D. Fedorets, L. Y. Gorelik, R. I. Shekhter, and M. Jonson, cond-mat/0311105 (unpublished).

[5] D. Fedorets, L. Y. Gorelik, R. I. Shekhter, and M. Jonson, cond-mat/0212561 (unpublished).

[6] T. Novotny, A. Donarini, and A.-P. Jauho, Phys. Rev. Lett. **90**, 256801 (2003).

[7] D. Fedorets, Phys. Rev. B **68**, 033106 (2003).

[8] A. Erbe, C. Weiss, W. Zwerger and R. H. Blick, Phys. Rev. Lett. **87**, 096106 (2001).

[9] D. V. Scheibe, A. Erbe and R. H. Blick, New J.Phys. **4**, 86 (2002).

[10] H. Park, J. Park, A. K. L. Lim, E. H. Anderson, A. P. Alivisatos, and P. L. McEuen, Nature **407**, 57 (2000).

[11] L. Y. Gorelik, R. I. Shekhter, V. M. Vinokur *et al.*, Phys. Rev. Lett. **91**, 088301 (2003).

[12] J. Park, A. Pasupathy, J. I. Goldsmith *et al.*, Nature **417**, 722 (2002).

[13] "Collosal Magnetoresistive Oxides" ed.by Y. Tokura (Gordon and Breach Science Publishers, Amsterdam, 2000).

[14] J. Z. Sun, W. J. Gallagher, P. R. Duncombe *et al.*, Appl. Phys. Lett. **69**, 3266 (1996).

[15] J. Z. Sun, L. Krusin-Elbaum, P. R. Duncombe *et al.*, Appl. Phys. Lett. **70**, 1769 (1997).

[16] Yu Lu *et al.*, Phys. Rev. B**54**, R8357 (1996).

[17] T. Kimura et al., Science, **274**, 1698 (1996).